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Peccei-Quinn field for inflation, baryogenesis, dark matter, and much more

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We propose a scenario of brane cosmology in which the Peccei-Quinn field plays the role of the inflaton and solves simultaneously many cosmological and phenomenological issues such as the generation of a heavy Majorana mass for the right-handed neutrinos needed for seesaw mechanism, MSSM μ-parameter, the right amount of baryon number asymmetry and dark matter relic density at the present universe, together with an axion solution to the strong CP problem without the domain wall obstacle. Interestingly, the scales of the soft SUSY-breaking mass parameter and that of the breaking of $U(1)_{\rm PQ}$ symmetry are lower bounded at $\mathcal{O}(10)$ TeV and $\mathcal{O}(10^{11})$ GeV, respectively.

INTRODUCTION

Probably, the most compelling source for the density perturbations in the present universe would be the classicalized quantum fluctuations of the inflaton [1, 2]. However, in addressing the spectral features of the density perturbations observed by Planck satellite [3], the flatness problem (called η-problem) of the inflaton potential [4, 5] turns out to be a big obstacle in the attempts of a simple realization of inflation based on an UV theory, e.g., supergravity or string theories with Einstein gravity.

Actually, the η-problem can be removed rather easily if the inflaton path is a simple compactified trajectory in a multi-dimensional field space without resorting to trans-Planckian excursions for the inflaton field [6–15] (see also [16]). Notably, the resulting inflation is effectively the same as single-field slow-roll inflation.

Another interesting possibility for tackling the η-problem is to introduce an extra-dimension [17]. Depending on the relative size of the energy density on the brane as compared with the brane tension, the expansion rate on the brane can be much larger than the one expected in Einstein gravity. This can resolve the η-problem, allowing inflation even with a steep potential [18].

In this letter, we show that the Peccei-Quinn field [19] of the DFSZ axion model [20, 21] in the supersymmetric framework is a natural candidate for the inflaton which can also trigger some necessary post-inflation cosmology such as baryogenesis and dark matter,when considering brane cosmology.

INFLATION ALONG A FLAT DIRECTION ON A BRANE

In Randall-Sundrum type II brane world scenario [17], after the bulk contribution to the brane is diluted away, Friedman equation on the brane can be written as [22]

$$3H^2 M_{\rm P}^2 = \rho \left(1 + \frac{\rho}{\Lambda}\right) \tag{1}$$

where Randall-Sundrum condition for vanishing cosmological constant is assumed, H is the expansion rate on the brane, $M_{\rm P} = 2.4 \times 10^{18}$ GeV is the reduced Planck mass, and ρ is the energy density on the brane other than the brane tension $\Lambda/2$ which is constrained to be larger than about $(2.3\,{\rm TeV})^4$ from the validity of the $(3+1)$-dimensional Newtonian gravity on scales larger than about 44μm [23, 24]. It implies that, if $\tilde{V} \equiv \rho/\Lambda \gg 1$ with ρ dominated by the potential energy V of a scalar field, a long enough inflation is possible even with a steep potential which would not allow or sustain an inflationary epoch in Einstein gravity. The slow-roll parameters of such an inflation on the brane are given by [25]

$$\epsilon \equiv -\frac{\dot{H}}{H^2} \approx \epsilon_E \frac{1+2\tilde{V}}{(1+\tilde{V})^2}, \tag{2}$$

$$\eta \equiv \frac{V''}{3H^2} \approx \frac{\eta_E}{1+\tilde{V}}, \tag{3}$$

$$\xi^2 \equiv \frac{V'V'''}{(3H^2)^2} \approx \frac{\xi_E^2}{(1+\tilde{V})^2} \tag{4}$$

where $\epsilon_{\rm E}, \eta_{\rm E}$ and $\xi_{\rm E}^2$ are the conventional slow-roll parameters in Einstein gravity. The spectral index and its running can be expressed as

$$\begin{aligned} n_s &= 1 - 6\epsilon + 2\eta, \\ \frac{dn_s}{d\ln k} &= 16\epsilon\eta - 24\epsilon^2 \frac{1+3\tilde{V}+3\tilde{V}^2}{(1+2\tilde{V})^2} - 2\xi^2 \end{aligned} \tag{5}$$

Notice that the spectral index is given by the same expression as in the case of Einstein gravity. The power spectrum and tensor-to-scalar ratio yield [25]

$$P_{\mathcal{R}} = \left(1+\tilde{V}\right)^3 P_{\mathcal{R},{\rm E}}, \quad r_T = \frac{16\epsilon}{1+2\tilde{V}} \tag{6}$$

with $P_{\mathcal{R},{\rm E}} = V/(24\pi^2 M_{\rm P}^4 \epsilon_{\rm E})$, while the number of e-foldings is

$$N_e = \int H dt = -\frac{1}{M_{\rm P}^2} \int \frac{V}{V'} \left(1+\tilde{V}\right) d\phi \tag{7}$$

with ϕ being the inflaton field. For the horizon scale of the present universe $k_H^{-1} \sim 6000$Mpc, the e-foldings expanded at the time the Planck pivot scale ($k_* =$

$0.05\mathrm{Mpc}^{-1}$) exits the horizon during inflation becomes [26]

$$
\begin{aligned}
N_{e,*} &\approx \frac{1}{3}\ln\left[\frac{\sqrt{6}}{16\pi}\mathbb{S}_0\left(\frac{k_0}{k_*}\frac{V_*}{V_e}\frac{V_e^{\frac{1}{4}}}{M_{\rm P}}\right)^3\left(\frac{T_{\rm d}}{\Lambda^{\frac{1}{4}}}\right)\tilde{V}_e^{\frac{5}{4}}\right] \\
&\simeq 39.4 - \ln\left[\frac{V_e}{V_*}\left(\frac{10^8\,{\rm GeV}}{V_e^{\frac{1}{4}}}\right)\right] - \frac{1}{3}\ln\left[\frac{\Lambda^{\frac{1}{4}}}{T_{\rm d}}\right] \\
&\quad + \frac{5}{12}\ln\tilde{V}_e
\end{aligned} \tag{8}
$$

where $\mathbb{S}_0$ is the total entropy of the present observable universe, the subscript '$_e$' in the right-hand side denotes a value at the end of inflation, and $T_{\rm d}$ is the decay temperature of the inflaton. Note that the last term in the second line of Eq. (8) is the new contribution coming from the enhancement of the expansion rate. Also, in order not to over-produce KK-gravitions, $T_{\rm d} \lesssim (\Lambda/2)^{1/4}$ is demanded [27].

Generically, thanks to the non-renormalization theorems, the potential V of a supersymmetric flat direction ($\Phi = \phi e^{i\theta}/\sqrt{2}$) is dominated by a mass term until the field gets close to the true vacuum or before it is lifted up by higher order terms. For a symmetry-breaking flat direction, the potential can be written as

$$
V = V_0 - \frac{1}{2}m^2\phi^2 + \ldots \tag{9}
$$

where V_0 is set for vanishing cosmological constant at the true vacuum, $m^2(>0)$ is a mass-square parameter, and '$\ldots$' represents higher order terms suppressed by a very large mass scale, e.g., Planck scale. Roughly, $V_0 \sim m^2\phi_0^2$ with ϕ_0 being the vacuum expectation value (VEV). Then, for $\phi \ll \phi_0$ one finds

$$
\frac{M_{\rm P}^2 V''}{V} \sim -\left(\frac{M_{\rm P}}{\phi_0}\right)^2, \quad \left|\frac{M_{\rm P}V'}{V}\right|^2 \sim \left(\frac{\phi}{\phi_0}\right)^2|\eta| \tag{10}
$$

Hence, if $\phi \ll \phi_0$ for a relevant cosmological scale during inflation, the slow-roll parameters of brane inflation are such that $\epsilon \ll \eta$ resulting in

$$
n_s - 1 \approx 2\eta \tag{11}
$$

Combining with the number of e-foldings in Eq. (7), we find

$$
N_{e,*}(n_s - 1) \sim 2\ln\left(\frac{\phi_*}{\phi_e}\right) \tag{12}
$$

Eq. (12) implies

$$
\phi_* \sim \phi_e e^{N_{e,*}(n_s-1)/2} \sim 0.37\phi_e = \mathcal{O}(0.1)\phi_0 \tag{13}
$$

where for a numerical estimation we used $N_{e,*} = 50$ and $\eta \approx -1.7\times10^{-2}$ [3] leading to

$$
\tilde{V} \sim 60\,(M_{\rm P}/\phi_0)^2 \tag{14}
$$

The power spectrum can be expressed as

$$
P_{\mathcal{R}} \sim \frac{1}{12\pi^2|\eta|^3}\left(\frac{m}{\phi_*}\right)^2 \tag{15}
$$

and $P_{\mathcal{R}} = 2.142\times10^{-9}$ which with $\eta = -1.7\times10^{-2}$ [3] leads to

$$
m/\phi_0 \sim 10^{-7} \tag{16}
$$

Also, from Eqs. (14) and (16), the scale of brane tension is found to be

$$
(\Lambda/2)^{1/4} \sim 0.6\left(\frac{m}{10^5\,{\rm GeV}}\right)^{1/2} m \tag{17}
$$

Since Λ is constrained as $(\Lambda/2)^{1/4} \gtrsim 2.3\,{\rm TeV}$, we find

$$
\mathcal{O}(10)\,{\rm TeV} \lesssim m \lesssim m_{\rm s} \tag{18}
$$

with $m_{\rm s}$ being the scale of the soft SUSY-breaking mass parameter, and

$$
\phi_0 \gtrsim \mathcal{O}(10^{11})\,{\rm GeV} \tag{19}
$$

It is very interesting to notice that, the Peccei-Quinn field is a natural and compelling candidate of ϕ.

THE MODEL

Motivated by the observation in the previous section, we consider the Peccei-Quinn field as the inflaton. For a specific realization, we consider a model defined by the following superpotential where gauge group and family indices are omitted:

$$
\begin{aligned}
W &= Y_u Q H_u \bar{u} + Y_d Q H_d \bar{d} + Y_e L H_d \bar{e} \\
&\quad + \frac{\lambda_N}{2} X N^2 + \lambda_\nu L H_u N \\
&\quad + \frac{\lambda_\mu X^2 H_u H_d}{M_{\rm P}} + \frac{\lambda_{XY} X^3 Y}{M_{\rm P}}
\end{aligned} \tag{20}
$$

here $U(1)$ Peccei-Quinn symmetry is assumed with charges assigned as $q_{\rm PQ}(X,Y) = (1,-3)$. The first line of Eq. (20) is the supersymmetric realization of the standard model (SM) Yukawa couplings. The second line triggers the seesaw mechanism when X develops a nonzero vacuum expectation value (VEV). That is, right-handed neutrinos (RHNs) get masses given by

$$
m_N = \lambda_N\langle X\rangle \tag{21}
$$

where the family index was suppressed. As $m_N \gg H$, RHNs are integrated out, replacing the second line of Eq. (20) by

$$
-\frac{1}{2}\frac{\lambda_\nu^2 (LH_u)^2}{\lambda_N\langle X\rangle} \tag{22}
$$

which generates Majora masses of active neutrinos at low energy given by

$$m_\nu = \frac{\lambda_\nu^2 v_u^2}{\lambda_N X_0} \tag{23}$$

with v_u being the VEV of H_u. The third line of Eq. (20) is nothing but a supersymmetric realization of the DFSZ axion model [28]. The first term of the third line reproduces the μ-term of the minimal supersymmetric standard model (MSSM) as X develops a non-zero VEV, and the second term stabilizes X and Y.

For a simple illustration of inflation and post-inflation cosmology in the model of Eq. (20), we set $X = Y = \Phi$ and consider a superpotential,

$$W = \frac{\lambda}{4}\frac{\Phi^4}{M_{\rm P}} \tag{24}$$

where $\lambda = \mathcal{O}(0.1-1)$ is a dimensionless constant. Then, including the soft SUSY-breaking terms with a negative mass-square parameter assumed [36], the scalar potential of Φ (Φ^*) becomes

$$V = V_0 - m^2|\Phi|^2 - \left(\frac{A\lambda}{4}\frac{\Phi^4}{M_{\rm P}} + {\rm c.c.}\right) + \left|\frac{\lambda\Phi^3}{M_{\rm P}}\right|^2 \tag{25}$$

where V_0 is again set to get a vanishing cosmological constant at true vacuum, and $m^2(>0)$ and A are soft SUSY-breaking parameters. The vacuum expectation value of Φ is found to be

$$\Phi_0 = \left(\frac{AM_{\rm P}}{6\lambda}\right)^{1/2}\left[1+\sqrt{1+\frac{12m^2}{A^2}}\right]^{1/2} \tag{26}$$

and

$$V_0 = \frac{2}{3}m^2|\Phi_0|^2\left[1+\frac{1}{24}\frac{A^2}{m^2}\left(1+\sqrt{1+\frac{12m^2}{A^2}}\right)\right] \tag{27}$$

The masses of radial and angular modes at true vacuum are

$$m_{\rm PQ,r}^2 = 2m^2 \tag{28}$$

$$m_{\rm PQ,a}^2 = \frac{2}{3}A^2\left[1+\sqrt{1+\frac{12m^2}{A^2}}\right] \tag{29}$$

COSMOLOGY

The model of Eq. (20) can accommodate inflation in a brane world scenario, and trigger the realization of a successful baryogenesis and supply dark matter production as follows.

Inflation

Although inflation in the model of Eq. (20) takes place along X, all its main features can be captured from the simplified example of Eq. (24). From our numerical analysis, we found that, as an example, the following set of parameters gives a perfect fit to the Planck data:

$$m \simeq 1.4\times10^5\,{\rm GeV}\,,\ \lambda\simeq 0.61,\ 2\tilde{V} = 10^{15.64} \tag{30}$$

which results in

$$m_{\rm PQ,r} = 1.96\times10^5\,{\rm GeV}\,,\ \phi_0 = 9.15\times10^{11}\,{\rm GeV} \tag{31}$$

and $(\Lambda/2)^{\frac{1}{4}} \simeq 51\,{\rm TeV}$. One can also lower ϕ_0 while keeping m/ϕ_0 nearly invariant for a correct value of $P_{\mathcal{R}}$. The decay rate of the flaton is assumed to be [32]

$$\Gamma_\phi = \frac{1}{4\pi}\left(1-\frac{B^2}{m_A^2}\right)^2\left(\frac{\mu}{m_{\rm PQ,r}}\right)^4\frac{m_{\rm PQ,r}^3}{\phi_0^2} \tag{32}$$

where B is the soft SUSY-breaking parameter of the Higgs bilinear term in MSSM and m_A is the mass of CP-odd Higgs. For $m_A = 1.03B$ and $\mu = m$, one finds $T_{\rm d} \simeq 550\,{\rm GeV}$ leading to $N_{e,*}\simeq 55$. Inflation ends at $\phi_e \simeq 0.8\phi_0$, and inflationary observables turn out to be

$$10^9 P_{\mathcal{R}} \simeq 2.136,\ n_s\simeq 0.962,\ \frac{dn_s}{d\ln k} = -6.1\times10^{-4} \tag{33}$$

The tensor-to-scalar ratio is miserably small, undistinguishable zero from an experimental point of view.

Even if Φ evolves in a 2-dimensional field space, for $\phi_* \lesssim \phi \lesssim \phi_e$, the angular direction of Φ is much flatter than radial direction. Hence, unless the trajectory of the inflaton is around the ridge of an angular potential (e.g., the potential involving A-parameter in Eq. (25)), the non-Gaussianities are expected to be $f_{\rm NL} \ll 1$, since perturbations are not enough to produce significant changes on the inflaton trajectories which nicely converge to the true vacuum. This is expected to be true as well for the real model of Eq. (20) where X plays the role of the slow-roll inflaton. Note that Y develops its VEV via the A-term of the second term in Eq. (20). So, Y is expected to follow its time-dependent vacuum position during the whole period of inflation, and therefore it would not affect the adiabatic density perturbations caused mainly by the perturbations of X.

A remark on the initial condition of the inflation along PQ-field is in order. If the PQ-field were the only light field, we can simply assume a radiation dominant universe before the onset of inflation. However, there are at least two possible candidates of light fields too: GUT Higgs and Planckian moduli. For $\Lambda^{1/4} = \mathcal{O}(10^{3-5})\,{\rm GeV}$ which is the case for the parameter set we have chosen for illustration purposes, they can also lead inflation while they roll out/in from/to the origin. As a simple possibility to realize PQ-inflation as the last inflation that may

be necessary to have a proper post-inflation cosmology (e.g., baryogenesis without unwanted relic problem), we may assume that the mass scales of those field are slightly higher than that of the PQ-field. In this case, the expansion rate at the end of the pre-existing inflation ($H_{\rm pre}$) could be a little bit larger than the one from PQ-inflation. The PQ-field would barely move from the origin once it is there. The potential danger of being drifted away by quantum fluctuations can be removed if PQ-field couples to a light degree of freedom around the origin. Specifically, in the presence of a coupling $W \supset \lambda_N \Phi N^2$, the effective mass-square of the PQ-field around the origin is

$$-m^2 + \lambda_N^2 \left(\frac{H_{\rm pre}}{2\pi} \right)^2 \tag{34}$$

since RHNs would have fluctuations of order $H_{\rm pre}/(2\pi)$. Hence, if

$$H_{\rm pre} > 2\pi m/\lambda_N \tag{35}$$

the PQ-field can be safely held at the origin. In other words, the mass scale of GUT Higgs or Planckian moduli is required to be larger than that of PQ-field by a factor about $2\pi/\lambda_N$. If GUT Higgs causes pre-inflation, the reheating temperature is likely to be low, since the interaction to light fields will be suppressed by GUT scale. In this case, the PQ-inflaton may start rolling out soon after the pre-inflation if $H_{\rm pre} \sim 2\pi m/\lambda_N$. On the other hand, Planckian moduli may end up at a symmetry-enhanced point leading to a rapid reheating and resulting in a rather high reheating temperature (e.g., $\sim \left(6 H_{\rm pre}^2 M_{\rm P}^2 \Lambda\right)^{1/8}$). In this case, the PQ-inflaton could be held at the origin until the temperature of thermal bath becomes comparable to m (a phase of thermal inflation [29, 30]).

Our choice of parameters results in $H \sim \mathcal{O}(10)\,(\Lambda/2)^{1/4}$ during inflation. In this case, the higher order kinetic terms of the inflaton, caused by brane fluctuations, are not suppressed enough, making the quadratic kinetic term a rather poor approximation [37]. It is non-trivial to properly handle this issue given the tachyonic nature of our inflaton, and out of the scope of this paper. However, at least, once the brane tension is given, enforcing the validity of considering only the quadratic kinetic term of the inflaton as we did in this work may impose a constraint on the energy density of the inflation V_0 such that Hubble fluctuation is at most of the same order as the brane tension. As a simple solution to this issue in our scenario, one can increase the size of A-parameter relative to the mass parameter m by a factor $\sim \mathcal{O}(10)$ in Eq. (25). In this case, the contribution of the A-term to the potential for the cosmological scales of interest becomes sizable, requiring a smaller ϕ_* in order to obtain the right amount of e-foldings. This means that ϵ_* becomes smaller. As a result, for a given brane tension one can lower the expansion rate during inflation by taking a slightly smaller mass scale for both parameters, m and A in Eq. (25), while keeping inflationary observables nearly unchanged. Baryogenesis and the production of dark matter discussed in the next subsections are barely affected although there will be minor changes in the working parameter space. So, we do not pursue the details of this case here.

Baryogenesis

The dynamical generation of μ can sustain a late time Affleck-Dine leptogenesis, as discussed in Refs. [31–33], if

$$m_{LH_u}^2 \equiv \left(m_L^2 + m_{H_u}^2\right)/2 < 0,\ 2m_{LH_u}^2 + |\mu|^2 > 0 \tag{36}$$

at scales below an intermediate scale, where m_L^2 and $m_{H_u}^2$ are respectively the soft SUSY-breaking mass-square parameters of L and H_u fields. In addition to Eq. (36), we also assume

$$|m_{LH_u}^2| \gtrsim m^2 \tag{37}$$

with $-m^2$ being understood as the soft SUSY-breaking mass-square of X so that the PQ-field rolls away from the origin following the LH_u-flat direction. Note that, in the model of Eq. (20) when X is held around the origin, the LH_u direction is stabilized by a quartic term. However, as X rolls out and becomes large, the renormalizable terms in the second line of Eq. (20) are replaced by the one of Eq. (22). The inverse of the time scale of the evolution of $\langle \ell \rangle$, the time-dependent VEV of the LH_u flat-direction, is

$$\langle \dot{\ell} \rangle / \langle \ell \rangle = \dot{X}/X = -\eta_X H \tag{38}$$

where η_X is the slow-roll parameter of the inflaton (X). It is much smaller than the mass scale of LH_u around the potential minimum. Hence, the LH_u flat-direction would trace its potential minimum as X evolves.

The lepton number asymmetry at the onset of the angular motion of the AD field is estimated to be

$$n_L \sim \frac{A_\nu}{m_{\rm PQ,r}} \frac{\lambda_\nu^2 \ell_e^4}{\lambda_N X_0} \sin(4\theta_\ell + \Delta\theta) \tag{39}$$

with $m_{\rm PQ,r}$ being the mass of the radial mode of X around the true vauum, and ℓ_e is the VEV of LH_u flat direction at the end of inflation. Because the oscillation of X causes time-dependent variations of the mass and of ℓ_e of LH_u, the initial lepton number asymmetry is suppressed by a factor which we denote as $\delta_{\rm s}$. Passing through the electroweak symmetry breaking, the asymmetry in the lepton number is converted to the one in baryon number via anomalous electroweak processes. Then, taking $n_B \sim n_L/3$, the late time baryon number

asymmetry after the inflaton decay can be expressed as

$$\frac{n_B}{s} \sim \delta_{\rm s}\frac{T_{\rm d}}{4m_{\rm PQ,r}}\left(\frac{A_\nu}{m_{\rm PQ,r}}\right)^2\left(\frac{\ell_0}{X_{\rm osc}}\right)^2 \times f(A_\nu)\left(\frac{\ell_e}{\ell_0}\right)^4\sin(4\theta_\ell+\Delta\theta) \tag{40}$$

where we use

$$\ell_0^2 = \frac{A_\nu M_\nu}{6\lambda_{\bar\nu}}\left[1+\sqrt{1+\frac{12|m_{LH_u}^2|}{A_\nu^2}}\right] \tag{41}$$

where $M_\nu \equiv \lambda_N X_0$ and $\lambda_{\bar\nu} \equiv \lambda_\nu^2$, $X_{\rm osc} = \alpha X_0$ is the initial oscillation amplitude of X at the end of inflation with $\alpha = \mathcal{O}(0.1)$, $f(A_\nu) \equiv \lambda_{\bar\nu}\ell_0^2/(A_\nu M_\nu)$, and $\Delta\theta$ is the relative CP-violating phase between the two lepton-number violating terms [33] and assumed to be of order one. The precise estimation of $\delta_{\rm s}$ requires a heavy numerical simulation which is out of the scope of this letter. However, from the simulation results of Ref. [32], we expect $\delta_{\rm s} = \mathcal{O}(10^{-3}-10^{-2})$, taking into account the larger expansion rate at the time of AD leptogenesis and the weaker preheating due to the smallness of the μ-term-induced mass variation of the LH_u direction. For the parameter set used in the previous section with $A_\nu \sim \sqrt{|m_{LH_u}^2|} \sim m$, we obtain $\ell_0 \sim 10^{9-10}\,{\rm GeV}$ and the factors appearing in the second line of Eq. (40) are $\mathcal{O}(0.1)$ (once taken together). Hence, the right amount of baryon number asymmetry would always be obtained if

$$T_{\rm d} = \mathcal{O}(10^{-3}-10^{-2})m \tag{42}$$

Notice that, unless $\delta_{\rm s}$ is too small, the right amount of baryon number asymmetry can be obtained by adjusting (mostly lowering down) $T_{\rm d}$ (or B/m_A - see Eq. (32)).

Since the AD field is a flat direction during inflation, there is a potential danger of generating too large baryon isocurvature perturbations ($\mathcal{S}_{b\gamma} \equiv \delta\ln(n_B/s)$), but they are naturally suppressed in our scenario for reasons which will be explained below.

Setting $\langle\theta_\ell\rangle = 0$ without loss of generality, from Eq. (39) one finds

$$\delta\ln\left(\frac{n_B}{s}\right) = -2\frac{\delta X_{\rm osc}}{X_{\rm osc}} + 4\left(\frac{\delta\ell_r}{\ell_0}+\cot(\Delta\theta)\frac{\delta\ell_a}{\ell_0}\right) \tag{43}$$

where $\delta\ell_{r,a}$ is the radial/angular perturbations of LH_u and we used $\delta\theta_\ell = \delta\ell_a/\ell_0$. Denoting the η of LH_u around the true vacuum as $\eta_{\ell,0}^{r,a}$ for the radial and angular modes of LH_u, the perturbations of each mode at the end of inflation can be written as

$$\delta\ell_{r,a} \sim \delta\ell_{r,a}^0 e^{-\eta_{\ell,0}^{r,a}N_{e,*}} \tag{44}$$

where $\delta\ell_{r,a}^0 = H/(2\pi)$. Note that $\eta_{\ell,0}^{r,a}$ can be larger than $|\eta_X|$ during inflation by a factor of a few, providing a suppression of $\delta\ell_{r,a}$. Observations constrain the baryon isocurvature perturbations to be $\beta_{\rm iso} \equiv P_{\mathcal{S}}/(P_{\mathcal{R}}+P_{\mathcal{S}}) \lesssim 0.038$ with $P_{\mathcal{S}}$ being the power spectrum of the isocurvature perturbations [34]. This can be interpreted as

$$|\mathcal{S}_{b\gamma}| \lesssim \frac{\Omega_{\rm DM}}{\Omega_B}\left(\frac{\beta_{\rm iso}}{1-\beta_{\rm iso}}P_{\mathcal{R}}\right)^{1/2} \lesssim 5.34\times10^{-5} \tag{45}$$

where $\Omega_{\rm DM}$ and Ω_B are the densities of cold dark matter and baryons, respectively. For $A_\nu = \sqrt{|m_{LH_u}^2|} = m$ as an example, $\mathcal{S}_{b\gamma} = 2.37\times10^{-5}(m/\ell_0/10^{-5})$. Therefore, we are safe as regards of baryon isocurvature perturbations.

Dark matter

If the inflaton decays well after the freeze-out of the lightest neutralino which is likely to be our case, in order to avoid dark matter over-production, the inflaton decays into flatinos and neutralinos should be kinematically forbidden (for example, see Ref. [35] for the parameter space). Hence, dark matter in our scenario is likely to be either neutralinos from thermal bath or axions, depending on the choice of parameters and the initial conditions of inflation.

The energy density stored in radiation after inflation until the energy density of PQ-field ($\rho_{\rm PQ}$) becomes comparable to the brane tension is expected to be

$$\rho_{\rm r} \sim \frac{\Gamma_{\rm d}}{H}\rho_{\rm PQ} \sim \sqrt{2}\left(\frac{\pi^2}{30}g_*(T_{\rm d})\right)^{1/2}T_{\rm d}^2\Lambda^{1/2} \tag{46}$$

where $g_*(T_{\rm d})$ is the number of relativistic degrees of freedom at inflaton decay. Hence, the background temperature during this epoch ($T_\times$) is nearly constant irrespective of H, and it may or may not be lower than the typical freeze-out temperature of neutralinos in Einstein gravity, $T_{\rm fz} \sim m_\chi/20$ with m_χ being the mass of neutralino.

If $T_\times > T_{\rm fz}$, for $T_{\rm d} < T_{\rm fz}$ the abundance of neutralinos would be diluted by the entropy production generated in the decay of the inflaton. The dilution factor is given by

$$\Delta = \left(\frac{T_{\rm d}}{\tilde T_{\rm d}}\right)^3 = 3^{2/3}\left(\frac{g_*(T_{\rm fz})}{g_*(T_{\rm d})}\right)^{2/3}\left(\frac{T_{\rm fz}}{T_{\rm d}}\right)^5 \tag{47}$$

where $\tilde T_{\rm d}$ is the would-be background temperature when $H \sim \Gamma_{\rm d}$ in case of no entropy production. The relic density of neutralino can then be expressed as

$$\frac{\Omega_{\rm DM}}{\Omega_{\rm DM}^{\rm obs}} \sim \left(\frac{m_\chi}{10^2\,{\rm GeV}}\right)^2\frac{1}{\Delta} \sim 10^{10}\left(\frac{m_\chi}{10^4\,{\rm GeV}}\right)^2\left(\frac{T_{\rm d}}{m_\chi}\right)^5 \tag{48}$$

where we used $g_*(T_{\rm fz}) = g_*(T_{\rm d})$ and $T_{\rm fz} = m_\chi/20$ for Δ. Hence, for $m_\chi \sim m_s \sim \mathcal{O}(10^{4-5})\,{\rm GeV}$ the abundance of neutralinos can match the observed relic density of dark matter if $T_{\rm d}/m_\chi \sim \mathcal{O}(10^{-3}-10^{-2})$ which may also be compatible with that required for baryogenesis.

If $T_\times < T_{\rm fz}$, neutralinos would never be in thermal equilibrium and their abundance will be negligible, its exact size depending on $T_\times / T_{\rm fz}$. In this case, dark matter would have to be mainly axions.

As PQ-symmetry is broken during inflation, the axion domain wall problem is absent, and the contribution to dark matter comes only from misalignment. In this case, the axion coupling constant can be much larger than the typically allowed upper-bound, since the misalignment angle is given as the initial condition of inflation.

CONCLUSIONS

In this letter, we proposed an inflation scenario along a supersymmetric flat direction in the context of Randall-Sundrum type-II brane world. Showing that, the mass scale (m) of the symmetry-breaking supersymmetric flat-direction is directly connected to the symmetry-breaking scale (ϕ_0) as $m/\phi_0 \sim 10^{-7}$, we identified the Peccei-Quinn field associated with the axion solution to the strong CP-prblem as the natural candidate to play the role of the inflaton. We also showed that a successful post-inflation cosmology including small active neutrino masses, baryogenesis and dark matter can be triggered by the non-zero vacuum expectation value of the PQ-inflaton via the generation of the heavy Majorna mass of the seesaw mechanism and the dynamical generation of the MSSM μ-parameter at the end of inflation. Either neutralinos or axions can be the main component of dark matter without an axion domain wall problem.

Our scenario does not have ad-hoc parameters except the brane tension (Λ). The fundamental Planck scale in five-dimensional spacetime is expected to be of $\mathcal{O}(10^9)$ GeV with a brane tension of $\mathcal{O}(10^{4-5})$ GeV. The brane tension imposes lower bounds on the scales of soft SUSY-breaking mass and breaking of $U(1)_{\rm PQ}$ symmetry at $\mathcal{O}(10)$ TeV and $\mathcal{O}(10^{11})$ GeV, respectively. This scenario can be excluded if primordial tensor perturbations are observed, a TeV scale SUSY (meaning $\Lambda^{1/4} < 1$ TeV) is found, or the axion coupling constant turns out to be lower than $\mathcal{O}(10^{10})$ GeV.

In summary our scenario is a full description of the early history of the Universe from inflation to Big-Bang nucleosynthesis that may be easily falsifiable at future collider experiments.

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